

Transverse Spin Distribution Function of Nucleon in Chiral Theory

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Abstract

At large N_c the nucleon can be viewed as a soliton of the effective chiral lagrangian. This picture of nucleons allows a consistent nonperturbative calculation of the leading-twist parton distributions at a low normalization point. We derive general formulae for the transverse spin quark distribution $h_1(x)$ in the chiral quark-soliton model. We present numerical estimates and compare them to the results obtained in other models.

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1. The transverse polarization distribution of quarks introduced by Ralston and Soper [1] and usually called $h_1(x)$ has gained an increasing amount of attention of theorists [2, 3, 4, 5, 6, 7]. The perspective of the experimental measurement of $h_1(x)$ [8, 9] has stimulated attempts to calculate $h_1(x)$ in various models of strong interactions [4, 5, 10, 11].

Recently a new approach to the calculation of quark distribution functions has been developed [12] within the context of the effective chiral quark-soliton model of nucleon [13]. In this paper we apply this approach to the calculation of the transverse spin distribution function $h_1(x)$. Although in reality the number of colours $N_c = 3$, the academic limit of large N_c is known to be a useful guideline. At large N_c the nucleon is heavy and can be viewed as a classical soliton of the pion field [14, 15]. In this paper we work with the effective chiral action given by the functional integral over quarks in the background pion field [16, 17, 18]:

$$\exp(iS_{\text{eff}}[\pi(x)]) = \int D\psi D\bar{\psi} \exp\left(i \int d^4x \bar{\psi}(i\rlap{\not{D}} - MU\gamma_5)\psi\right),$$

$$U = \exp[i\pi^a(x)\tau^a], \quad U\gamma_5 = \exp[i\pi^a(x)\tau^a\gamma_5] = \frac{1+\gamma_5}{2}U + \frac{1-\gamma_5}{2}U^\dagger. \quad (1)$$

Here ψ is the quark field, M is the effective quark mass which is due to the spontaneous breakdown of chiral symmetry and U is the $SU(2)$ chiral pion field. The effective chiral action given by eq. (1) is known to contain automatically the Wess–Zumino term and the four-derivative Gasser–Leutwyler terms, with correct coefficients. Eq. (1) has been derived from the instanton model of the QCD vacuum [18, 19], which provides a natural mechanism of chiral symmetry breaking and enables one to express the dynamical mass M and the ultraviolet cutoff intrinsic in eq. (1) through the Λ_{QCD} parameter. It should be mentioned that eq. (1) is of a general nature: one need not believe in instantons and still can use eq. (1).

An immediate application of the effective chiral theory (1) is the quark-soliton model of baryons of ref. [13], which is in the spirit of the earlier works [20, 21]. According to this model nucleons can be viewed as $N_c (=3)$ “valence” quarks bound by a self-consistent pion field (the “soliton”) whose energy coincides in fact with the aggregate energy of the quarks of the negative-energy Dirac continuum. Similarly to the Skyrme model large N_c is needed as a parameter to justify the use of the mean-field approximation, however, the $1/N_c$ corrections can be and in some cases have been computed [22, 23, 24, 26, 27]. The quark-soliton model of nucleons developed in ref. [13] includes a collective-quantization procedure to deal with the rotational excitations of the quark-pion soliton.

Turning to the calculation of the transverse spin distribution function $h_1(x)$ we note that the model possesses all features needed for a successful description of the nucleon parton structure: it is an essentially quantum field-theoretical relativistic model with explicit quark degrees of freedom, which allows an unambiguous identification of quark as well as antiquark distributions in the nucleon. This should be contrasted to the Skyrme model where it is not too clear how to define quark and antiquark distributions.

2. The transverse spin quark distribution function for flavour f is defined as follows [4, 5]

$$h_1^f(x) = \int \frac{d\lambda}{4\pi} e^{i\lambda x} \langle P, S_T | \bar{\psi}_f(0) \rlap{\not{\epsilon}} \gamma_5 \rlap{\not{S}}_T \psi_f(\lambda n) | P, S_T \rangle. \quad (2)$$

Below we work in the nucleon rest frame where the transverse nucleon spin S_T and the light-cone vector n can be chosen as follows:

$$n^\mu = \frac{1}{M_N}(1, -1, 0, 0), \quad S_T^\mu = (0, 0, 0, 1). \quad (3)$$

The operators in eq. (2) depend on the QCD normalization point μ . In contrast to QCD, the effective chiral field theory is nonrenormalizable and contains an explicit ultraviolet cut-off. In the instanton vacuum model this cutoff appears as the inverse average instanton size: $\bar{\rho}^{-1} \approx 600$ MeV [19, 25]. The results obtained below thus refer to a low point normalization point μ of the order of 600 MeV.

Now we can rewrite (2) in the form

$$h_1^f(x) = -\frac{1}{4\pi} \int_{-\infty}^{\infty} dz^0 e^{ixM_N z^0} \times \langle P, S_T | \psi_f^+(0)(1 + \gamma^0 \gamma^1) \gamma_5 \gamma^3 \psi_f(z) | P, S_T \rangle \Big|_{z^1=-z^0, z^2=z^3=0}. \quad (4)$$

Let us remind the reader how the nucleon is described in the effective low-energy theory (1). Integrating out the quarks in (1) one finds the effective chiral action,

$$S_{\text{eff}}[\pi^a(x)] = -N_c \text{Sp} \log D(U), \quad D(U) = i\partial_0 - H(U), \quad (5)$$

where $H(U)$ is the one-particle Dirac hamiltonian,

$$H(U) = -i\gamma^0 \gamma^k \partial_k + M\gamma^0 U^{\gamma_5}, \quad (6)$$

and $\text{Sp} \dots$ denotes a functional trace.

For a given time-independent pion field $U = \exp(i\pi^a(\mathbf{x})\tau^a)$ one can determine the spectrum of the Dirac hamiltonian,

$$H\Phi_n = E_n\Phi_n. \quad (7)$$

It contains the upper and lower Dirac continua (distorted by the presence of the external pion field), and, in principle, also discrete bound-state level(s), if the pion field is strong enough. If the pion field has unity winding number, there is exactly one bound-state level which travels all the way from the upper to the lower Dirac continuum as one increases the spatial size of the pion field from zero to infinity [13]. We denote the energy of the discrete level as E_{lev} , $-M \leq E_{\text{lev}} \leq M$. One has to occupy this level to get a non-zero baryon number state. Since the pion field is colour blind, one can put N_c quarks on that level in the antisymmetric state in colour.

The limit of large N_c allows us to use the mean-field approximation to find the nucleon mass. To get the nucleon mass one has to add $N_c E_{\text{lev}}$ and the energy of the pion field. Since the effective chiral lagrangian is given by the determinant (5) the energy of the pion field coincides exactly with the aggregate energy of the lower Dirac continuum, the free continuum subtracted. The self-consistent pion field is thus found from the minimization of the functional [13]

$$M_N = \min_U N_c \left\{ E_{\text{lev}}[U] + \sum_{E_n < 0} (E_n[U] - E_n^{(0)}) \right\}. \quad (8)$$

From symmetry considerations one looks for the minimum in a hedgehog ansatz:

$$U_c(\mathbf{x}) = \exp[i\pi^a(\mathbf{x})\tau^a] = \exp[in^a\tau^a P(r)], \quad r = |\mathbf{x}|, \quad \mathbf{n} = \frac{\mathbf{x}}{r}, \quad (9)$$

where $P(r)$ is called the profile of the soliton.

The minimum of the energy (8) is degenerate with respect to translations of the soliton in space and to rotations of the soliton field in ordinary and isospin space. For the hedgehog field (9) the two rotations are equivalent. Quantizing slow rotations of the saddle-point pion field [15, 13] leads to the projection on a nucleon with given spin (S_3) and isospin (T_3) components which includes the integration over all spin-isospin $SU(2)$ orientation matrices R of the soliton,

$$\langle S = T, S_3, T_3 | \dots | S = T, S_3, T_3 \rangle = \int dR \phi_{S_3 T_3}^{S=T}(R) \dots \phi_{S_3 T_3}^{S=T}(R). \quad (10)$$

Here $\phi_{S_3 T_3}^{S=T}(R)$ is the rotational wave function of the nucleon given by the Wigner finite-rotation matrix [15, 13]:

$$\phi_{S_3 T_3}^{S=T}(R) = \sqrt{2S+1}(-1)^{T+T_3} D_{-T_3, S_3}^{S=T}(R). \quad (11)$$

3. Eq. (2) contains a matrix element of a non-local quark bilinear operator in the nucleon state with definite 4-momentum P and spin and isospin components. According to [12] one can write a general expression for such matrix elements; the time dependence of the quark operators is accounted for by the energy exponents. Taking nucleon at rest, we can write this matrix element as a sum over all occupied quark states

$$\begin{aligned} \langle S_3, T_3 | \psi_{fi}^\dagger(x^0, \mathbf{x}) \psi^{gj}(y^0, \mathbf{y}) | S_3, T_3 \rangle &= 2M_N N_c \int d^3\mathbf{X} \int dR \phi_{S_3 T_3}^*(R) \\ &\times \sum_{\substack{n \\ \text{occup.}}} \exp[iE_n(x^0 - y^0)] \Phi_{n, f'i}^\dagger(\mathbf{x} - \mathbf{X}) (R^\dagger)_f^{f'} R_{g'}^g \Phi_n^{g'j}(\mathbf{y} - \mathbf{X}) \phi_{S_3 T_3}(R). \end{aligned} \quad (12)$$

Here we have written explicitly all the flavour ($f, g = 1, 2$) and the Dirac ($i, j = 1, \dots, 4$) indices for clarity. The functions Φ_n are eigenstates of energy E_n of the Dirac hamiltonian (6) in the external (self-consistent) pion field U_c . Summation over colour indices is implied in the quark bilinears, hence the factor N_c .

Applying the general formula of the effective chiral theory (12) to the calculation of the transverse spin distribution function (4) we obtain in the leading order of the $1/N_c$ expansion

$$\begin{aligned} h_1^f(x) &= -\frac{N_c M_N}{2\pi} \int dR |\phi_{T_3 S_3}(R)|^2 \int_{-\infty}^{\infty} dz^0 e^{ixM_N z^0} \sum_{\substack{n \\ \text{occup.}}} e^{-iE_n z^0} \\ &\times \langle n | R^\dagger T^f R (1 + \gamma^0 \gamma^1) \gamma_5 \gamma^3 \exp(-iz^0 p^1) | n \rangle. \end{aligned} \quad (13)$$

Here T^f is the matrix projecting onto the flavour f . We use notation $p^k = -i\partial_k$ for the operator of the 3-momentum. One can easily check that in the leading order of the $1/N_c$ expansion only the *isovector* polarized distribution survives:

$$h_1^u(x) - h_1^d(x) = -\frac{N_c M_N}{2\pi} \int dR |\phi_{T^3 S^3}(R)|^2 \int_{-\infty}^{\infty} dz^0 e^{ix M_N z^0} \sum_{\substack{n \\ \text{occup.}}} e^{-i E_n z^0} \\ \times \langle n | R^\dagger \tau^3 R (1 + \gamma^0 \gamma^1) \gamma_5 \gamma^3 \exp(-iz^0 p^1) | n \rangle. \quad (14)$$

The integral over the orientation matrix R can be easily computed. As a result we obtain

$$h_1^u(x) - h_1^d(x) = \frac{N_c M_N}{6\pi} \int_{-\infty}^{\infty} dz^0 e^{ix M_N z^0} \sum_{\substack{n \\ \text{occup.}}} e^{-i E_n z^0} \\ \times \langle n | \tau^3 (1 + \gamma^0 \gamma^1) \gamma_5 \gamma^3 \exp(-iz^0 p^1) | n \rangle. \quad (15)$$

Taking into account the hedgehog symmetry (9) we can replace here $1 \leftrightarrow 3$. Passing to the momentum representation and integrating over z^0 , we obtain the expression convenient for practical calculations:

$$h_1^u(x) - h_1^d(x) = \frac{N_c M_N}{6\pi} \int \frac{d^2 p_\perp}{(2\pi)^2} \sum_{\substack{n \\ \text{occup.}}} \\ \times \langle n | \mathbf{p} \rangle (1 + \gamma^0 \gamma^3) \gamma_5 \gamma^1 \tau^1 \langle \mathbf{p} | n \rangle \Big|_{p^3 = x M_N - E_n}. \quad (16)$$

It is well known that the transverse spin distribution function h_1 satisfies the sum rule [4, 5]

$$\int_{-1}^1 dx h_1^f(x) = g_T^f \quad (17)$$

where g_T^f is the tensor charge of the nucleon defined as follows

$$\langle P, S | \bar{\psi}_f [\gamma_\mu, \gamma_\nu] \psi_f | P, S \rangle = g_T^f \bar{u}(P, S) [\gamma_\mu, \gamma_\nu] u(P, S). \quad (18)$$

Here $u(P, S)$ is the standard Dirac spinor describing the nucleon with momentum P and spin S .

Substituting expression for $h_1(x)$ (16) into the lhs of the sum rule (17) and extending the integral over x to the whole real axis (which is justified by the large N_c limit [12]) we obtain

$$\int_{-1}^1 dx h_1^f(x) = \frac{N_c}{3} \sum_{\substack{n \\ \text{occup.}}} \langle n | \gamma_5 \gamma^1 \tau^1 | n \rangle. \quad (19)$$

The result in the rhs coincides with the expression obtained in paper [26] for the tensor charge of nucleon so that the sum rule (17) automatically holds in the effective chiral theory.

4. In eq. (16) one has to sum over all occupied levels of the Dirac Hamiltonian, including both negative Dirac continuum and discrete level.

In contrast to the longitudinal spin quark distribution function $g_1(x)$ whose Dirac continuum contribution is ultraviolet divergent [12], in the transverse case the continuum contribution is ultraviolet finite. This can be seen from the gradient expansion for $h_1(x)$ where the first nonvanishing term appears in higher orders compared to the gradient expansion of $g_1(x)$. This also hints that the continuum contribution to $h_1(x)$ is relatively small, whereas the continuum contribution to the longitudinal spin distribution function $g_1(x)$ is rather sizable [12]. An additional argument in favour of the suppression of the continuum contribution to h_1 comes from the numerical calculation of the tensor charge performed in [26, 27], which shows that the contribution of the Dirac continuum is essentially smaller than the contribution of the discrete level. Since the sum rule (17) holds in our model, it is natural to expect that the discrete level gives the dominant contribution also to the quark distribution function $h_1(x)$.

The bound-state level occurs [13] in the grand spin $K = 0$ and parity $\Pi = +$ sector of the Dirac hamiltonian (6). In that sector the eigenvalue equation takes the form

$$\begin{pmatrix} M \cos P(r) & -\frac{\partial}{\partial r} - \frac{2}{r} - M \sin P(r) \\ \frac{\partial}{\partial r} - M \sin P(r) & -M \cos P(r) \end{pmatrix} \begin{pmatrix} h_0(r) \\ j_1(r) \end{pmatrix} = E_{\text{lev}} \begin{pmatrix} h_0(r) \\ j_1(r) \end{pmatrix}. \quad (20)$$

We assume that the radial wave functions are normalized by the condition

$$\int_0^\infty dr r^2 [h_0^2(r) + j_1^2(r)] = 1. \quad (21)$$

Introducing the radial wave functions in the momentum representation

$$h(k) = \int_0^\infty dr r^2 h_0(r) \sqrt{\frac{k}{r}} J_{\frac{1}{2}}(kr), \quad j(k) = \int_0^\infty dr r^2 j_1(r) \sqrt{\frac{k}{r}} J_{\frac{3}{2}}(kr), \quad (22)$$

we can write the expression for the level contribution to h_1 in the form

$$\begin{aligned} [h_1^u(x) - h_1^d(x)]_{\text{lev}} &= \frac{1}{3} N_c M_N \int_{|xM_N - E_{\text{lev}}|}^\infty \frac{dk}{2k} \left\{ h^2(k) + \frac{(xM_N - E_{\text{lev}})^2}{k^2} j^2(k) \right. \\ &\quad \left. - 2 \frac{(xM_N - E_{\text{lev}})}{k} h(k) j(k) \right\}. \end{aligned} \quad (23)$$

5. Recently the following inequality has been suggested by Soffer [28]

$$q^f(x) + g_1^f(x) \geq 2|h_1^f(x)|. \quad (24)$$

Its status in QCD has been studied in [29]. Let us check this inequality in the chiral quark soliton model. A delicate point is that in the leading order of the $1/N_c$ expansion the unpolarized quark distribution q^f is saturated by the singlet part whereas g_1^f and h_1^f get the

main contribution from their nonsinglet parts. Therefore in the large N_c approach the above inequality can be rewritten as follows:

$$[q^u(x) + q^d(x)] - |g_1^u(x) - g_1^d(x)| \geq 2|h_1^u(x) - h_1^d(x)|. \quad (25)$$

Note that the structure of the expressions for all quark distribution function in the effective quark soliton model entering inequality (25) is essentially the same. They contain a sum of diagonal matrix elements over occupied states of the Dirac hamiltonian, differing only by the spin and isospin matrices Γ appearing in these matrix elements:

$$\begin{aligned} [q^u(x) + q^d(x)], \quad \Gamma_q &= (1 + \gamma^0 \gamma^3), \\ [g_1^u(x) - g_1^d(x)], \quad \Gamma_L &= -\frac{1}{3}(1 + \gamma^0 \gamma^3) \gamma_5 \tau^3, \\ [h_1^u(x) - h_1^d(x)], \quad \Gamma_T &= \frac{1}{3}(1 + \gamma^0 \gamma^3) \gamma_5 \gamma^1 \tau^1 \end{aligned} \quad (26)$$

(the explicit expressions for $q^f(x)$ and $g_1(x)$ can be found in ref. [12]). Taking into account the following inequalities (understood in the matrix sense)

$$|(1 + \gamma^0 \gamma^3) \gamma_5 \tau^3| \leq 1 + \gamma^0 \gamma^3, \quad (27)$$

$$|(1 + \gamma^0 \gamma^3) \gamma_5 \gamma^1 \tau^1| \leq 1 + \gamma^0 \gamma^3 \quad (28)$$

one can easily see that the Soffer inequality (24) holds for the contribution of each separate occupied level.

One should keep in mind that in our model the quark distribution functions $q(x)$ and $g_1(x)$ are ultraviolet divergent and in principle the regularization can violate the above argument. However, at least in the limit of large cutoff we can prove the validity of the Soffer inequality. Indeed, function $h_1(x)$ is ultraviolet finite whereas $q(x)$ and $g_1(x)$ are ultraviolet divergent. The analysis of these divergences has been performed in ref. [12] and it follows from this analysis that the ultraviolet divergence of the combination $[q^u(x) + q^d(x)] - |g_1^u(x) - g_1^d(x)|$ is given by an explicitly positive expression so that in the limit of large cutoff this expression definitely becomes larger than the ultraviolet finite $h_1(x)$, which proves that the Soffer inequality holds in this limit.

6. We have calculated numerically the level contribution to the transverse isovector quark distribution. In this calculation we use the parameters obtained by using a variational estimate of the profile of the pion-field soliton (see eq. (9)) performed in ref. [13] yielding for $M = 350$ MeV

$$P(r) = -2 \arctan \left(\frac{r_0^2}{r^2} \right), \quad r_0 \approx 1.0/M, \quad M_N \approx 1170 \text{ MeV}. \quad (29)$$

This profile function has a correct behaviour at small and large distances and is stable with respect to small perturbations. In our numerics we use the analytic profile (29).

The results of our calculations for g_1 and h_1 are presented at Fig.1. The antiquark distributions are shown at Fig.2. From the latter plot we can see that the antiquark transverse spin distribution is essentially smaller than the corresponding longitudinal distribution which

is related to the fact that the longitudinal distribution gets a large contribution from the quarks of the Dirac continuum [12] in contrast to the transverse distribution.

Comparing our results to the calculations in other models – bag model [4, 5], QCD sum rules [10], chiral chromodielectric model [11], one can see that the main source of difference is the $1/N_c$ approximation used in our calculations. As it was mentioned above, in the leading order of the $1/N_c$ expansion we have $h_1^u(x) = -h_1^d(x)$. This relation is much more general than the model we used and also holds in the large N_c QCD. It is well known that typically the large N_c counting is a good guide line for understanding nonperturbative physics. In order to understand the reliability of the $1/N_c$ expansion in the case of h_1 it were instructive to compute the next-order $1/N_c$ corrections. Unfortunately, the calculation of these corrections in our model is rather complicated in the case of the quark distribution functions. However, these $1/N_c$ corrections have been computed for the tensor charge [26, 27] and this calculation shows that these corrections really enhance g_T^u and suppress g_T^d . This means the $1/N_c$ corrections of our model can reconcile the large N_c approach with the model calculations of [4, 5, 10, 11].

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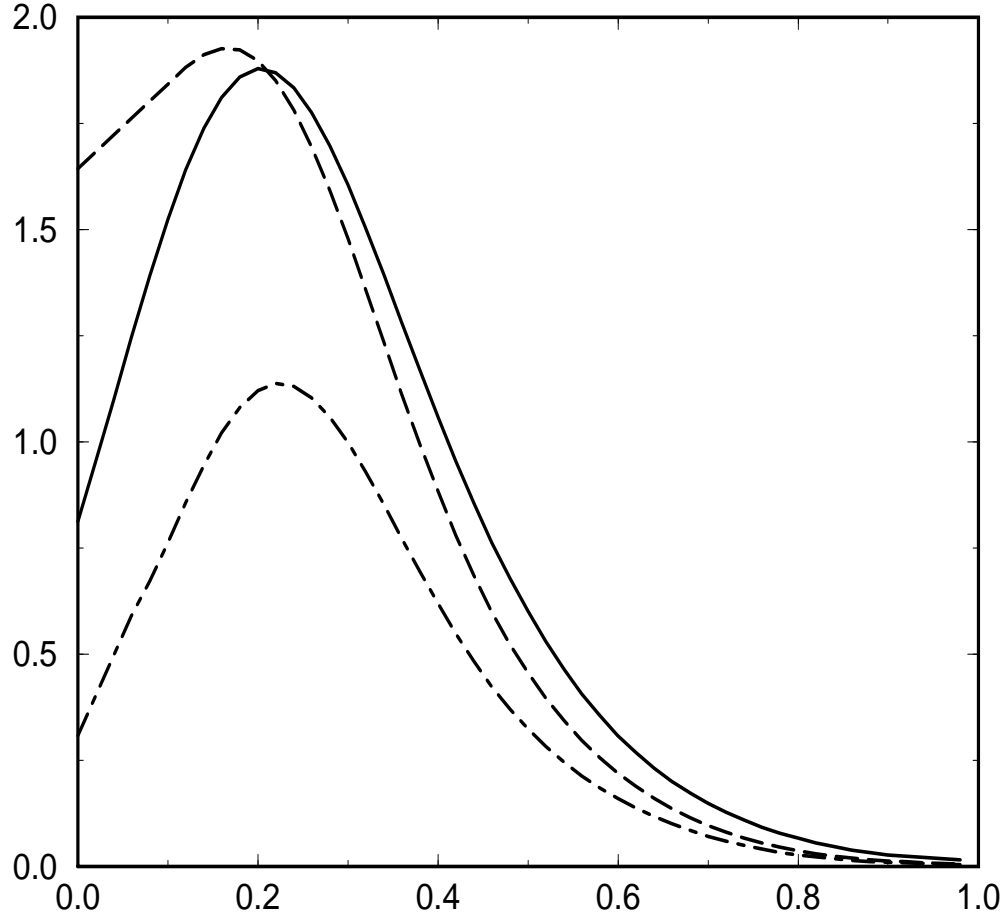


Figure 1: The transverse spin quark distribution function $h_1^u - h_1^d$ (solid line) plotted against the longitudinal quark distribution $g_1^u - g_1^d$ (dashed) computed in the same model. The contribution of the discrete level to $g_1^u - g_1^d$ is depicted by the dot-dashed line.

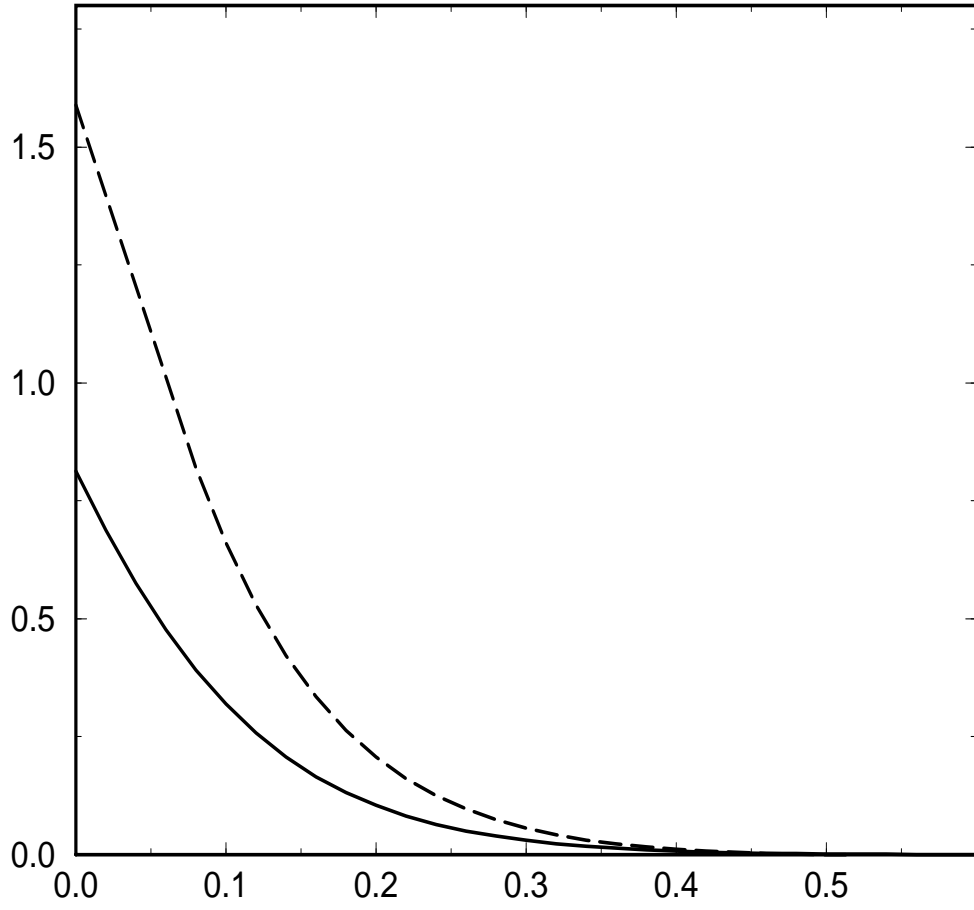


Figure 2: The transverse spin antiquark distribution function $\bar{h}_1^d - \bar{h}_1^u$ (solid line) plotted against the longitudinal antiquark distribution $\bar{g}_1^u - \bar{g}_1^d$ (dashed) computed in the same model.